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# Riemann tensor of the ambient universe, the dilaton, and Newton's constant

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We investigate a four-dimensional world, embedded into a five-dimensional spacetime, and find the five-dimensional Riemann tensor via generalization of the Gauss(-Codacci) equations. We then derive the generalized equations of the four-dimensional world and also show that the square of the dilaton field is equal to Newton's constant. We find plausible constant and nonconstant solutions for the dilaton.

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Since the pioneering work of Kaluza [1] and Klein [2], who unified gravitation with electromagnetism, the implications of possible extra dimensions to our world have been under intense investigation—see [3] for an extensive collection of papers on higher-dimensional unification. Jordan and Müller [4] and Thiry [5] used the equations of Kaluza-Klein theory to show that the gravitational constant can be expressed as a dynamical field. A constant solution for Newton's constant, however, is then allowed only if the square of the Maxwell electromagnetic tensor vanishes. Here we examine a *dual* setup in which this problem can be avoided. Based on the original Kaluza model, we are here considering a general embedding of a four-dimensional world into a five-dimensional ambient spacetime. We derive a generalization of the Gauss(-Codacci) equations by utilizing all degrees of freedom (the entire geometry) of the ambient spacetime, and we also show how they affect the physics of the four-dimensional world. As a result we find a system of equations for the electromagnetic field, gravitational field, and dilaton field. One of these equations is a plausible generalization of the gauge fixed Maxwell equations in the presence of a dilaton field. We also show that the square of the dilaton field is equal to (modulo numerical factors) the Newton's constant. The gauge freedom of the electromagnetic fields is transferred to a freedom in fixing the dilaton field. Apart from the constant solution for the dilaton, we give an example of a nonconstant solution describing a time-varying Newton constant in an expanding universe (see also [6–11] and others). We also give a general formula for generating different solutions for the dilaton field.

We consider a four-dimensional world  $M$ , embedded into a five-dimensional spacetime  $V$  (see [12–22] and references therein for a detailed discussion on embeddings). Let  $y^i$  ( $i = 1, 2, 3, 4$ ) denote the coordinates on  $M$  and  $x^\mu$  ( $\mu = 1, 2, 3, 4, 5$ ) denote the coordinates on  $V$ . Greek indices will be related to the five-dimensional spacetime  $V$ , while latin indices will be associated with the four-dimensional spacetime  $M$ . Let  $\psi(x^\mu) = s = \text{const}$  be the equation of the hyper-

surface  $M$ . We assume that this equation is given *a priori*. The physical principles motivating the choice of type of hypersurface [specified by the function  $\psi(x^\mu)$ ] and the choice of a particular four-dimensional spacetime slice (given by  $s$ ) are open problems for multidimensional theories and not the subject of our analysis. We will retain  $s$  as a parameter of the model.

One can alternatively express the parametric equations of  $M$  as  $x^\mu = x^\mu(y^j, s)$  and treat the parameter  $s$  as a coordinate; this then represents a coordinate transformation with inverse

$$y^j = y^j(x^\mu), \quad s = \psi(x^\mu). \quad (1)$$

We assume that this transformation is invertible at each point. This means that the Jacobi matrices of the transformation and its inverse have nonvanishing determinants everywhere. Thus, to globally parametrize the foliated five-dimensional spacetime  $V$ , it is sufficient to use the coordinates of the four-dimensional world  $M$  and the foliation parameter  $s$ .

The vector normal to the surface is

$$N_\mu = \frac{\partial \psi}{\partial x^\mu}. \quad (2)$$

Let us also define

$$e_j^\mu = \frac{\partial x^\mu}{\partial y^j}, \quad n^\mu = \frac{\partial x^\mu}{\partial s}, \quad E_\mu^i = \frac{\partial y^i}{\partial x^\mu}. \quad (3)$$

The derivatives are related as follows:

$$\partial_k \equiv \frac{\partial}{\partial y^k} = e_k^\mu \partial_\mu, \quad (4)$$

$$\partial_\nu \equiv \frac{\partial}{\partial x^\nu} = E_\nu^k \partial_k + N_\nu \partial_s. \quad (5)$$

Obviously, if we denote  $e_5^\mu = n^\mu$  and  $E_\mu^5 = N_\mu$ , then  $(e_\nu^\mu)$  and  $(E_\nu^\mu)$  will be the Jacobi matrices of the transformation  $(x^\mu) \rightarrow (y^i, s)$  and its inverse. Therefore,

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$$e_\sigma^\mu E_\nu^\sigma = \delta_\nu^\mu = E_\sigma^\mu e_\nu^\sigma. \quad (6)$$

This orthogonality condition is equivalent to

$$e_i^\mu E_\mu^j = \delta_i^j, \quad (7)$$

$$N_\mu n^\mu = 1, \quad (8)$$

$$E_\mu^i e_i^\nu + N_\mu n^\nu = \delta_\mu^\nu, \quad (9)$$

$$N_\mu e_j^\mu = 0, \quad (10)$$

$$n^\mu E_\mu^j = 0. \quad (11)$$

(Note that  $\delta_\mu^\mu = \dim V = 5$  and  $\delta_i^i = \dim M = 4$ .)

Thus the bases  $(e_j^\mu, n^\mu)$  and  $(E_\mu^j, N_\mu)$  are dual. They do not depend on the metric of either spacetime, but only on the particular embedding chosen.

Let us now introduce a scalar field  $\phi(y^k, s)$ , a vector field  $A_i(y^k, s)$ , and the metric tensor  $g_{ij}(y^k, s)$ , which are  $M$ -valued tensor functions on  $V$ . We further define the metric  $G_{\mu\nu}$  of the five-dimensional spacetime  $V$  as an expansion over the basis vectors  $E_\mu^i$  and  $N_\mu$ :

$$G_{\mu\nu} = E_\mu^i E_\nu^j g_{ij} + (N_\mu E_\nu^i + N_\nu E_\mu^i) A_i + N_\mu N_\nu \phi. \quad (12)$$

Taking  $x^i = y^i$ ,  $x^5 = s = \text{const}$ , i.e.,  $e_i^\mu = \delta_i^\mu$ ,  $n^\mu = \delta_5^\mu$ ,  $E_\mu^i = \delta_\mu^i$ ,  $N_\mu = \delta_\mu^5$  in Eq. (12) corresponds to the original Kaluza model [1]. Klein's modification [2]  $g_{ij} \rightarrow g_{ij} + A_i A_j$ , together with the identification of  $\phi$  as a dilaton is the model put forward by Jordan and Müller [4] and Thiry [5]. We note that the metric (12) has the same form as the inverse of the metric of Klein's model, and thus the two theories are dual: Kaluza's model corresponds to slicing, while Klein's model corresponds to threading of the five-dimensional spacetime [15]. The case with  $A_i = 0$  has also been considered (see, for example, [22,23] and references therein).

The lack of gauge invariance for the fields  $A_i$ , which we nevertheless will associate with the electromagnetic potentials, in view of the slicing-threading duality, is compensated by the freedom to fix  $\phi$ . This, as will become clear later, is the freedom to fix the dilaton field.

Returning to Eq. (12), we define  $g^{ij}$  as the inverse of the metric  $g_{ij}$ . Thus  $A^i = g^{ij} A_j$  and  $A^2 = g^{ij} A_i A_j$ .

The inverse  $G^{\mu\nu}$  of the metric  $G_{\mu\nu}$  on  $V$  is then given by

$$G^{\mu\nu} = h^{ij} e_i^\mu e_j^\nu - \theta A^i (e_i^\mu n^\nu + e_i^\nu n^\mu) + \theta n^\mu n^\nu, \quad (13)$$

where  $\theta = (\phi - A^2)^{-1}$  and  $h^{ij} = g^{ij} + \theta A^i A^j$ . One can easily check that  $G^{\mu\lambda} G_{\lambda\nu} = \delta_\nu^\mu$ . Using the inverse  $G^{\mu\nu}$  (13) of the metric  $G_{\mu\nu}$ , we can raise and lower the five-dimensional indices to get

$$N^\mu = G^{\mu\nu} N_\nu = \theta (n^\mu - A^i e_i^\mu), \quad (14)$$

$$N^2 = N_\mu N^\mu = \theta, \quad (15)$$

$$n_\mu = G_{\mu\nu} n^\nu = A_i E_\mu^i + N_\mu \phi, \quad (16)$$

$$n^2 = n_\mu n^\mu = \phi. \quad (17)$$

Note that when  $A^i = 0$  and  $\phi = 1$ , then  $N^\mu = n^\mu$  as in the Arnowitt-Deser-Misner approach [22].

The extrinsic curvature is

$$K_{jl} = e_j^\mu e_l^\nu Q_{\mu\nu}, \quad (18)$$

where  $Q_{\mu\nu} = \nabla_\mu N_\nu = \nabla_\nu N_\mu$ .

Multiplying Eq. (18) across by  $E_\alpha^j E_\beta^l$  and applying the orthogonality conditions (7)–(11), one easily finds

$$Q_{\alpha\beta} = E_\alpha^i E_\beta^j K_{ij} + (N_\alpha E_\beta^i + N_\beta E_\alpha^i) f_i + N_\alpha N_\beta \chi, \quad (19)$$

with  $\chi = n^\mu n^\nu Q_{\mu\nu}$  and

$$f_j = n^\mu e_j^\beta Q_{\mu\beta} = A^i K_{ij} + \frac{1}{N} \partial_j N, \quad (20)$$

where  $N = \sqrt{N_\mu N^\mu}$ .

The four-dimensional Christoffel symbols

$$\gamma_{jk}^i = \frac{1}{2} g^{il} (\partial_k g_{lj} + \partial_j g_{lk} - \partial_l g_{jk}) \quad (21)$$

can then be expressed as

$$\gamma_{jl}^i = (\partial_j e_l^\mu + e_j^\alpha e_l^\beta \Gamma_{\alpha\beta}^\mu) E_\mu^i - K_{jl} A^i, \quad (22)$$

where  $\Gamma_{\alpha\beta}^\mu$  are the five-dimensional Christoffel symbols. Further, the four-dimensional Riemann curvature tensor

$$r_{jkl}^i = \partial_k \gamma_{jl}^i - \partial_l \gamma_{jk}^i + \gamma_{jl}^m \gamma_{mk}^i - \gamma_{jk}^m \gamma_{ml}^i, \quad (23)$$

using Eqs. (22) and (7)–(10), becomes

$$\begin{aligned} r_{jkl}^i &= E_\alpha^i e_j^\lambda e_k^\mu e_l^\nu R_{\lambda\mu\nu}^\alpha + E_\alpha^i (\nabla_\mu n^\alpha) (e_j^\mu K_{lk} - e_k^\mu K_{lj}) \\ &\quad + \nabla_l (A^i K_{kj}) - \nabla_k (A^i K_{lj}) + A^i A^m (K_{ml} K_{kj} - K_{mk} K_{lj}), \end{aligned} \quad (24)$$

where  $R_{\lambda\mu\nu}^\alpha$  is the five-dimensional Riemann curvature tensor. The above is a *generalization* of the Gauss equations. Taking  $A^i = 0$  and  $\phi = 1$ , one simply recovers the well known Gauss equations (see, for example, [24])

$$r_{jkl}^i = E_\alpha^i e_j^\lambda e_k^\mu e_l^\nu R_{\lambda\mu\nu}^\alpha + K_k^i K_{lj} - K_l^i K_{kj}. \quad (25)$$

Let us now expand the five-dimensional Riemann curvature tensor over our basis. Using its symmetries we can write

$$\begin{aligned} R_{\lambda\mu\nu\sigma} &= E_\lambda^i E_\mu^j E_\nu^k E_\sigma^l U_{ijkl} + [(N_\lambda E_\mu^j - N_\mu E_\lambda^j) E_\nu^k E_\sigma^l \\ &\quad + (N_\nu E_\sigma^j - N_\sigma E_\nu^j) E_\lambda^k E_\mu^l] V_{jkl} \\ &\quad + (N_\lambda E_\mu^j - N_\mu E_\lambda^j) (N_\nu E_\sigma^l - N_\sigma E_\nu^l) W_{jl}, \end{aligned} \quad (26)$$

where the coefficients in this expansion satisfy

$$U_{ijkl} = U_{klji} = -U_{ijlk} = -U_{jilk}, \quad (27)$$

$$V_{jkl} = -V_{jlk}, \quad W_{jl} = W_{lj}. \quad (28)$$

Using Eq. (24), one can further find

$$U_{ijkl} = e_i^\lambda e_j^\mu e_k^\nu e_l^\sigma R_{\lambda\mu\nu\sigma} \\ = r_{ijkl} - (\pi_{ik}\pi_{lj} - \pi_{il}\pi_{kj}), \quad (29)$$

$$V_{jkl} = n^\lambda e_j^\mu e_k^\nu e_l^\sigma R_{\lambda\mu\nu\sigma} \\ = A^i U_{ijkl} - \frac{1}{N} (\nabla_k \pi_{lj} - \nabla_l \pi_{kj}), \quad (30)$$

where  $\pi_{jl} = (1/N)K_{jl}$ .

Finding the remaining tensor  $W_{jl}$  is more complicated. One can easily see that

$$W_{jl} = n^\mu n^\sigma e_j^\lambda e_l^\nu R_{\lambda\mu\nu\sigma} = S_{jl} + A^k V_{lkj}, \quad (31)$$

where

$$S_{jl} = \frac{1}{N^2} n^\sigma e_j^\lambda e_l^\nu (\nabla_\lambda Q_{\sigma\lambda} - \nabla_\sigma Q_{\nu\lambda}). \quad (32)$$

In the above we identify derivatives in the direction of  $n^\sigma$ . To handle this type of term, we will have to explicitly invoke the dependence on  $s$ .

First, by virtue of Eq. (5), we get the following expression for the extrinsic curvature (18):

$$K_{jl} = -\frac{N^2}{2} (\nabla_j A_l + \nabla_l A_j) + \frac{N^2}{2} \partial_s g_{jl} \\ - \frac{1}{2} N^\lambda (g_{kl} \partial_j E_\lambda^k + g_{kj} \partial_l E_\lambda^k + A_l \partial_j N_\lambda + A_j \partial_l N_\lambda) \\ + \frac{N^2}{2} G_{\mu\nu} [A^k \partial_k (e_j^\mu e_l^\nu) - \partial_s (e_j^\mu e_l^\nu)]. \quad (33)$$

Using Eqs. (20), (31), (32), and (33), it follows that

$$W_{jl} = \frac{1}{N^3} \nabla_j \nabla_l N - \frac{2}{N^4} (\partial_j N) (\partial_l N) + A^i A^k U_{ijkl} \\ + \frac{1}{N} [\nabla_j (A^k \pi_{kl}) + \nabla_l (A^k \pi_{kj})] - \frac{1}{N} A^k \nabla_k \pi_{jl} \\ + \frac{1}{N^2} \pi_{kl} \pi_j^k - \frac{1}{N} \partial_s \pi_{jl} + \Omega_{jl}, \quad (34)$$

where  $\Omega_{jl}$  contains only terms that are proportional to derivatives of the basis vectors and their duals with respect to  $(y^k, s)$ .

The five-dimensional Ricci tensor can easily be calculated from Eq. (26):

$$R_{\mu\nu} = E_\mu^j E_\nu^l [h^{ik} U_{ijkl} - N^2 A^k (V_{jkl} + V_{lkj}) + N^2 W_{jl}] + (N_\mu E_\nu^l \\ + N_\nu E_\mu^l) (-h^{ik} V_{jkl} + N^2 A^j W_{jl}) + N_\mu N_\nu h^{jl} W_{jl}. \quad (35)$$

Then the five-dimensional Einstein equations in vacuum

$$R_{\mu\nu} = 0 \quad (36)$$

reduce to

$$h^{ik} U_{ijkl} - N^2 A^k (V_{jkl} + V_{lkj}) + N^2 W_{jl} = 0, \quad (37)$$

$$h^{jk} V_{jkl} - N^2 A^j W_{jl} = 0, \quad (38)$$

$$h^{jl} W_{jl} = 0. \quad (39)$$

Multiplying Eq. (37) by  $A^j$  and adding it to Eq. (38) allows us to exclude  $W_{jl}$  from Eq. (38). Then, using the expressions (29) and (30) for  $U_{ijkl}$  and  $V_{jkl}$ , Eq. (38) becomes

$$\nabla_k \pi_l^k - \nabla_l \pi_k^k = 0. \quad (40)$$

Equations (40) (as we will see below) are a *generalization* of Maxwell's equations in a fixed gauge.

One has to make a very important point here. Klein's theory corresponds to a threading decomposition of the five-dimensional spacetime [15]. Rigorous analysis [16] shows that the curvature tensor of the hypersurface formed is given by Zelmanov's curvature tensor, which differs from the ordinary Riemann curvature tensor by additional terms containing  $s$  derivatives of the four-dimensional metric [17–21]. Independence of the parameter  $s$  forces the two curvature tensors to be equal and thus represents a surface forming condition. In Kaluza's theory, the foliation of the five-dimensional spacetime corresponds to slicing [15]. Then the four-dimensional metric  $g_{ij}$  naturally appears as the slicing metric, and imposing independence on the parameter  $s$  is not at all necessary.

To simplify the analysis of the physics described by the fields  $A_i$ ,  $\phi$ , and  $N$ , we will, however, put aside the  $s$ -dependent terms. Also, for simplicity, we will assume that the basis elements and their duals are constant (thus recovering the original Kaluza theory). The tensor  $\Omega_{jl}$  will then vanish from Eq. (34).

Equation (40) becomes

$$\nabla_k F^{kl} = -2A_k r^{kl} + \frac{2}{N^2} (\pi^{kl} - \pi_j^j g^{kl}) \partial_k N. \quad (41)$$

Here  $F_{kl} = \nabla_k A_l - \nabla_l A_k$  is the Maxwell electromagnetic tensor with  $A_k$  being the electromagnetic potential.

The first term on the right-hand side of Eq. (41) describes an interaction between electromagnetic and gravitational fields. We assume that it is much smaller than the remaining terms, so that we can neglect it. Note that  $A_k$  cannot be “gauged up” to increase the scale of  $A_k r^{kl}$ . Furthermore, if  $N$  is a constant, then Eq. (41) becomes the usual Maxwell equations

$$\nabla_k F^{kl} = 0. \quad (42)$$

The remaining two equations are

$$\nabla_k \left( \frac{f^k}{N} \right) = 0, \quad (43)$$

$$r_{jl} - \frac{1}{2} g_{jl} r = \frac{N^2}{2} T_{jl}, \quad (44)$$

where  $r = g^{ik} r_{ik}$  and  $r_{jl}$  are the four-dimensional scalar curvature and four-dimensional Ricci tensor, respectively.

The energy-momentum tensor  $T_{jl}$  is therefore given by

$$T_{jl} = T_{jl}^{\text{Maxwell}} + g^{ik} \nabla_i B_{jlk} + C_{jl} + D_{jl}, \quad (45)$$

where

$$T_{jl}^{\text{Maxwell}} = g^{ik} F_{ij} F_{kl} - \frac{1}{4} g_{jl} F_{ik} F^{ik},$$

$$B_{jlk} = A_k \nabla_l A_j - A_l \nabla_k A_j - A_j F_{kl} + \nabla_j (A_k A_l) + g_{jl} (A^i \nabla_k A_i - A_k \nabla^i A_i), \quad (46)$$

$$C_{jl} = g_{jl} A^i A^k r_{ik} - 2A^i A_l r_{ij} - 2A^i A_j r_{il}, \quad (47)$$

$$D_{jl} = \frac{4}{N^4} (\partial_j N) (\partial_l N) - \frac{2}{N^3} \nabla_j \nabla_l N - \frac{2}{N^2} \pi_k^k (A_l \partial_j N + A_j \partial_l N) + \frac{2}{N^2} [-A^k \pi_{jl} + A_l \pi_j^k + A_j \pi_l^k - g_{jl} (A^i \pi_i^k - A^k \pi_i^i)] \partial_k N. \quad (48)$$

We will analyze each of these terms separately. The first one  $T_{jl}^{\text{Maxwell}}$  is the Maxwell energy-momentum tensor. The tensor  $C_{jl}$  describes the interaction between the electromagnetic and gravitational fields. From Eq. (44) we see that, if  $N^2$  is very small (as we will confirm later), then  $r_{jl}$  will be of the order of  $N^2$ , which justifies the neglect of the interaction terms in Eq. (41) and the tensor  $C_{jl}$ .

Using Eq. (33) in Eq. (20) and then Eq. (20) in Eq. (43), we see that a constant solution for  $N$  is allowed by Eq. (43) if  $\phi$  satisfies

$$\nabla^k \partial_k \phi = \frac{1}{2} F^{ik} F_{ik}, \quad (49)$$

where  $F^{ik}$  is a solution of Eq. (42). For the constant solution for  $N$ , the tensor  $D_{jl}$  vanishes. Moreover,

$$0 \equiv g^{ij} \nabla_i T_{jl} = g^{ij} \nabla_i T_{jl}^{\text{Maxwell}}, \quad (50)$$

since  $g^{ml} g^{nk} \nabla_m \nabla_n B_{jlk} = -(2/N) \nabla_j \nabla_k (f^k/N) = 0$  in view of Eq. (43).

In other words, the conservation law (50) is given by the usual Maxwell energy-momentum tensor  $T_{jl}^{\text{Maxwell}}$  and  $N^2$  plays the role of the Newton constant  $G_N$ :

$$\frac{N^2}{2} = \frac{8\pi G_N}{c^4}. \quad (51)$$

The generalized Einstein and Maxwell equations (41)–(48) will be modified further upon inclusion of  $s$ -derivative terms. Equation (51), however, holds regardless of whether the  $s$ -derivative terms are included or not. One has to point out here that in the setup of Thiry [5], and in [14],  $G_N \approx \phi^2$ , where  $\phi$  satisfies Eq. (49). This implies that a constant solution for  $\phi$  and, consequently,  $G_N$  is possible only when the unphysical constraint  $F^{ik} F_{ik} = 0$  is satisfied.

In contrast, in the dual setup, a constant solution is possible. However,  $N$  (together with  $A_i$  and  $g_{ij}$ ) is a solution to

the system of equations (41), (43), and (44) and, in general, does not need to be a constant. Then it plays the role of a dilaton field.

To illustrate this, consider the standard cosmological metric [11] with  $E_\mu^\nu = \delta_\mu^\nu$ :

$$ds_{(5)}^2 = -s^2 dt^2 + t^{2/\alpha} s^{2/(1-\alpha)} (dr^2 + r^2 d\Omega^2) + \alpha^2 (1-\alpha)^{-2} t^2 ds^2. \quad (52)$$

Changing variables by  $r \rightarrow s^\gamma e^{\beta r}$  with  $\gamma = -(1/2)(1+\alpha)/(1-\alpha)$  and  $t^{1/\alpha} \rightarrow a(t)$ , we get

$$ds_{(5)}^2 = -s^2 \alpha^2 [a(t)]^{2\alpha-2} \dot{a}^2(t) dt^2 + 2A_r dr ds + s a^2(t) e^{2\beta r} (\beta^2 dr^2 + d\Omega^2) + \phi ds^2, \quad (53)$$

where  $A_r = \gamma \beta a^2(t) e^{2\beta r}$ ,  $A_t = A_\varphi = A_\theta = 0$ , and  $\phi = (\gamma^2/s) a^2(t) e^{2\beta r} + \alpha^2 (1-\alpha)^{-2} a^2(t)$ .

We take  $\beta$  to be a negative constant, so that the field  $A_r$  will fall off toward infinity. Since  $a(t)$  describes the inflation, we note that the field  $A_r$  expands as  $a^2(t)$ . The dilaton (which models Newton's constant) varies as

$$N^2 = (1-\alpha)^2 \alpha^{-2} [a(t)]^{-2\alpha}. \quad (54)$$

Thus,

$$\frac{\dot{G}}{G} = -2\alpha \frac{\dot{a}}{a} = -2\alpha H, \quad (55)$$

where  $H$  is Hubble's constant. Observational limits [7] put  $\alpha < 10^{-3}$ .

One should note that the four-dimensional metric is now  $s$  dependent, but this does not pose a problem in the slicing formulation. Only the term  $\partial_s g_{jl}$  from the extrinsic curvature (33) should be recovered.

Finally, we note that the general solution for the dilaton field can be written as

$$N^2 = \frac{(\det g_{ik})(\det E_\nu^\mu)^2}{\det G_{\mu\nu}}, \quad (56)$$

for a solution  $G_{\mu\nu}$  of Eq. (36) and embedding specified with  $E_\nu^\mu$ .

To recapitulate, we have found plausible generalizations of the Einstein-Maxwell equations and explained the origin of the constant solution for the dilaton (representing Newton's constant  $G_N$ ) as well as the possibilities for modeling nonconstant solutions for different cosmologies (representing time-varying  $G_N$ ) in relation to the gauge freedom of our model.

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